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Description of $\rho(1700)$ with the FCA method

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We study the $\rho K\bar{K}$ system with an aim to describe the $\rho(1700)$ resonance. The chiral unitary approach has achieved success in a description of systems of the light hadron sector. With this method, the $K\bar{K}$ system in the isospin sector $I=0$, is found to be a dominant component of the $f_0(980)$ resonance. Therefore, by regarding the $K\bar{K}$ system as a cluster, the $f_0(980)$ resonance, we evaluate the $\rho K\bar{K}$ system applying the fixed center approximation to the Faddeev equations. We construct the ρK unitarized amplitude using the chiral unitary approach. As a result, we find a peak in the three-body amplitude around 1739 MeV and a width of about 227 MeV. The effect of the width of ρ and $f_0(980)$ is also discussed. We associate this peak to the $\rho(1700)$ which has a mass of 1720 ± 20 MeV and a width of 250 ± 100 MeV.

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We study the $\rho K\bar{K}$ system in the sector $I(J^P)=1(1^-)$ within the fixed center approximation (FCA) to obtain the $\rho(1700)$ resonance. A pair of $K\bar{K}$ is assumed to form the scalar cluster, the $f_0(980)$ resonance because the $K\bar{K}$ component in $f_0(980)$ is found to be dominant [1]. In order to obtain the unitarized ρK amplitude, we follow the schemes given by refs. [2, 3] and extend them to the isospin $I=3/2$ sector. Basically we follow the formalism given by refs. [4, 5, 6, 7].

To implement the Faddeev equation within the fixed center approximation, we need the ρK ($\rho\bar{K}$) unitarized amplitude. Here we utilize the amplitude given in the previous work [2, 3] as to the vector-pseudoscalar interaction in the sector with strangeness $S=1$ and isospin $I=1/2$. Following the Bethe-Salpeter approach, we have the VP two-body scattering amplitude as

$$T=[1+V\hat{G}]^{-1}(-V)\vec{\varepsilon}\cdot\vec{\varepsilon}', \tag{1}$$

where V is an interaction kernel which will be discussed later, $\hat{G}$ is $(1+\frac{1}{3}\frac{q_l^2}{M_l^2})G$ being a diagonal matrix and $\vec{\varepsilon}(\vec{\varepsilon}')$ represents a polarization vector of the incoming (outgoing) vector-meson. Thanks to the on-shell factorization and the dimensional regularization, a loop function of pseudoscalar and vector mesons G_l can be expressed as a function of the energy $\sqrt{s}$

$$\begin{aligned} G_l(\sqrt{s}) \;=\; & \frac{1}{16\pi^2}\left\{a(\mu)+\ln\frac{M_l^2}{\mu^2}+\frac{m_l^2-M_l^2+s}{2s}\ln\frac{m_l^2}{M_l^2}\right. \\ & +\frac{q_l}{\sqrt{s}}\left[\ln(s-(M_l^2-m_l^2)+2q_l\sqrt{s})+\ln(s+(M_l^2-m_l^2)+2q_l\sqrt{s})\right. \\ & \left.\left.-\ln(-s+(M_l^2-m_l^2)+2q_l\sqrt{s})\;-\ln(-s-(M_l^2-m_l^2)+2q_l\sqrt{s})\right]\right\}, \end{aligned} \tag{2}$$

where a momentum q_l is determined at the center of mass frame and μ is a scale parameter in this scheme. Furthermore in accordance with ref. [3], we take into account the effect of the propagation of unstable particles in terms of the Lehmann representation. Instead of the original loop function eq. (2), we use

$$\tilde{G}_l(\sqrt{s}) \;=\; \frac{1}{C_l}\int_{(M_l-2\Gamma_l)^2}^{(M_l+2\Gamma_l)^2} ds_V G_l(\sqrt{s},\sqrt{s_V},m_l)\times\left(-\frac{1}{\pi}\right)\mathrm{Im}\left\{\frac{1}{s_V-M_l^2+iM_l\Gamma_l}\right\}, \tag{3}$$

with the normalization for the lth component

$$C_l \;=\; \int_{(M_l-2\Gamma_l)^2}^{(M_l+2\Gamma_l)^2} ds_V\times\left(-\frac{1}{\pi}\right)\mathrm{Im}\left\{\frac{1}{s_V-M_l^2+iM_l\Gamma_l}\right\}, \tag{4}$$

with m_l, M_l, Γ_l, the mass of the pseudoscalar meson, mass of the vector and width of the vector respectively. We can obtain the interaction kernel V by the use of the WCCWZ approach [8, 9, 10] where the interaction Lagrangian stems from a nonlinear realization of chiral symmetry. By projecting over s-wave, we have the VP potential

$$V_{ij}(s) \;=\; -\frac{\vec{\varepsilon}\cdot\vec{\varepsilon}'}{8f^2}C_{ij}\left[3s-(M_i^2+m_i^2+M_j^2+m_j^2)-\frac{1}{s}(M_i^2-m_i^2)(M_j^2-m_j^2)\right], \tag{5}$$

where the index $i(j)$ represents the VP channel of the incoming (outgoing) particles. The coefficients C_{ij} in eq. (5) for the $I=1/2$ and $I=3/2$ sector are tabulated in tables 1 and 2 respectively.

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	ϕK	ωK	ρK	$K^*\eta$	$K^*\pi$
ϕK	0	0	0	$-\sqrt{\frac{3}{2}}$	$-\sqrt{\frac{3}{2}}$
ωK	0	0	0	$\frac{\sqrt{3}}{2}$	$\frac{\sqrt{3}}{2}$
ρK	0	0	-2	$-\frac{3}{2}$	$\frac{1}{2}$
$K^*\eta$	$-\sqrt{\frac{3}{2}}$	$\frac{\sqrt{3}}{2}$	$-\frac{3}{2}$	0	0
$K^*\pi$	$-\sqrt{\frac{3}{2}}$	$\frac{\sqrt{3}}{2}$	$\frac{1}{2}$	0	-2

Table 1: Coefficients C_{ij} in the $I = 1/2$ sector.

	ρK	$K^*\pi$
ρK	1	1
$K^*\pi$	1	1

Table 2: Coefficients C_{ij} in the $I = 3/2$ sector.

In order to have an appropriate unitarized amplitude, we use the following parameter set chosen to reproduce $K_1(1270)$ in ref. [3] as $\mu = 900$ MeV, $a(\mu) = -1.85$, $f = 115$ MeV .

Under the fixed center approximation [4, 5, 6, 7], it is assumed that K and $\bar{K}$ cluster together and the structure of the cluster is kept against the collision of ρ. This idea leads the three-body scattering amplitude T which can read a summation of the two following partition functions T_1 and T_2

$$\begin{aligned} T_1 &= t_1 + t_1 G_0 T_2, \\ T_2 &= t_2 + t_2 G_0 T_1, \\ T &= T_1 + T_2, \end{aligned} \tag{6}$$

where the subscripts $i = 1(2)$ of T_i and t_i represent the component particle $K(\bar{K})$ in the cluster and the diagrammatic sketches are depicted in fig. 1. In the present work, we can rewrite $t = t_1 = t_2$ and t is given as a mixture of the different isospin states of $I = 1/2$ and $I = 3/2$, $t = \left(2t_{\rho K}^{I=3/2} + t_{\rho K}^{I=1/2}\right)/3$, where $t_{\rho K}$ is the ρK unitarized scattering amplitude given by eq. (1). By adopting the field normalization used in refs. [5, 6, 7], G_0 of eq. (6) reads as a function of the energy $\sqrt{s}$

$$G_0(\sqrt{s}) = \frac{1}{2M_{f_0}} \int \frac{d^3q}{(2\pi)^3} F_{f_0}(q) \frac{1}{q^{02}(s) - \vec{q}^{\,2} - m_\rho^2 + i m_\rho \Gamma_\rho}, \tag{7}$$

where M_{f_0} is the mass of the $f_0(980)$ resonance and the width of the ρ is taken into account in the above propagator. The energy of the propagator q^0 is determined at the three-body rest frame, $q^0(\sqrt{s}) = (s + m_\rho^2 - M_{f_0}^2)/2\sqrt{s}$. Here we utilize the form factor F_{f_0} to give the momentum distribution of f_0 to the G_0 function by the use of the formalism developed in refs. [11, 12, 13]

$$F_{f_0}(q) = \frac{1}{\mathcal{N}} \int_{\substack{p<k_{\max} \\ |\vec{p}-\vec{q}|<k_{\max}}} d^3p \left(\frac{1}{2\omega_K(\vec{p})}\right)^2 \frac{1}{M_{f_0} - 2\omega_K(\vec{p})}$$

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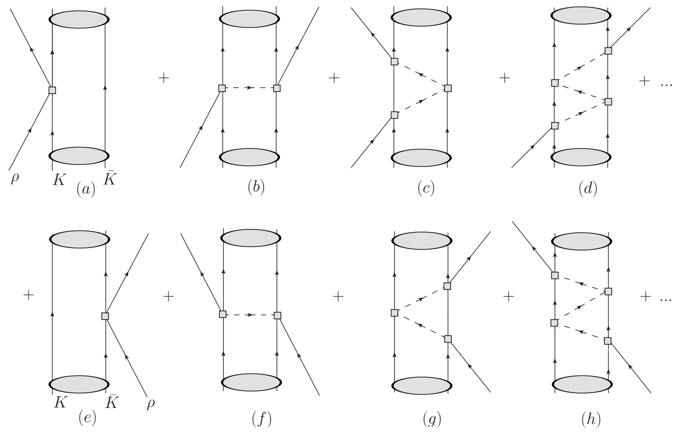

Figure 1: Diagrammatic representation of the fixed center approximation for the $\rho K\bar{K}$ system.

$$\times \left(\frac{1}{2\omega_K(\vec{p}-\vec{q})}\right)^2 \frac{1}{M_{f_0}-2\omega_K(\vec{p}-\vec{q})}, \tag{8}$$

where the normalization $\mathscr{N}$ is given by

$$\mathscr{N} = \int_{p<k_{\max}} d^3p \left[\left(\frac{1}{2\omega_K(\vec{p})}\right)^2 \frac{1}{M_{f_0}-2\omega_K(\vec{p})}\right]^2. \tag{9}$$

From ref. [1], we take $k_{\max} = \sqrt{\Lambda^2 - m_K^2}$ and $\Lambda = 1030$ MeV for getting the $f_0(980)$ from the $K\bar{K}$ cluster. Considering the S-matrix, we have a simple expression of the three-body scattering amplitude

$$T(s) = 2\left[\tilde{t}(s') + \tilde{t}(s')G_0(s)\tilde{t}(s') + \cdots\right] = 2\frac{\tilde{t}(s')}{1-\tilde{t}(s')G_0(s)}, \tag{10}$$

where t is replaced by $\tilde{t} = (2m_{f_0}/2m_K)t$ coming from the normalization of the fields and note that the unitarized amplitude t is a function of $s' = 1/2(s + M_\rho^2 + 2m_K^2 - M_{f_0}^2)$. Finally we consider the width of f_0 by replacing the mass of the cluster M_{f_0} in eqs. (8) and (9) with $M_{f_0} - i\Gamma_{f_0}/2$. The amplitude with the $f_0(980)$ and ρ width effect is shown in fig. 2 and the masses and widths of the dynamically generated state are listed in table 3. It is shown that the inclusion of the $f_0(980)$ width induces a suppression of the magnitude of the peak and the peak becomes broader as the width of the $f_0(980)$ increases. Furthermore it is also a remarkable feature that the peak position is not so affected by this prescription.

Through this work, we construct the $\rho K\bar{K}$ three-body amplitude by means of the fixed center approximation. In our framework, a pair of $K\bar{K}$ is considered to form a scalar meson cluster $f_0(980)$, based on ref. [1]. We use the ρK unitarized amplitude provided by refs. [2, 3] in a manner giving a respect to chiral symmetry. In the three-body amplitude, we have a peak at the energy around 1748 MeV rather independent of the width of the $f_0(980)$. Besides, it is seen that the inclusion of the $f_0(980)$ width makes the peak wider and gives a good agreement with the

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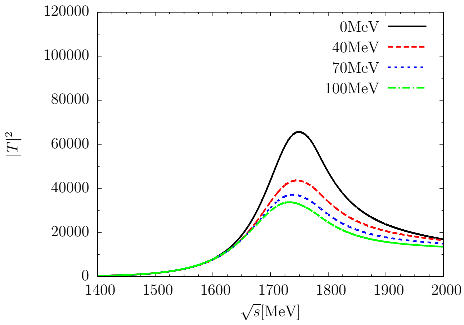

Figure 2: The $\rho K\bar{K}$ amplitude with the width effect for $\Gamma_{f_0} = 0, 40, 70, 100$ MeV, respectively.

	$\Gamma_{f_0} = 0$	$\Gamma_{f_0} = 40$	$\Gamma_{f_0} = 70$	$\Gamma_{f_0} = 100$	PDG [14]
Mass	1748.0	1743.6	1739.2	1734.8	1720 $\pm$20
Width	160.8	216.4	227.2	224.6	250 $\pm$100

Table 3: The masses and widths of dynamically generated states with the width effects. (in MeV)

experimental data of the $\rho(1700)$, both for the position and the width. Since the ρ decays into $\pi\pi$ mostly, the above results might be related to the dominant decay mode of the $\rho(1700)$, $\rho\pi\pi$ and 4π. Our approach to the $\rho K\bar{K}$ system provides the description of the $\rho(1700)$ as a dynamically generated state and then we conclude that the building block of the $\rho(1700)$ resonance are the ρ and $f_0(980)$.

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